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Light Stops in the MSSM: Implications for Photino Dark Matter and Top Quark Decay

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Abstract

We consider the viability of the minimal supersymmetric standard model with a light ($m_{\tilde{t}_1} <$ 45 GeV) stop. In order for its relic abundance to be cosmologically significant, the photino as dark matter must be quite close in mass to the stop, $(m_{\tilde{t}_1} - m_{\tilde{\gamma}}) \simeq 3 - 7$ GeV. However, as we show, the photino despite its low mass is virtually undetectable by either direct or indirect means. We also discuss the implications of these masses on the top quark branching ratios.

Despite its name, the minimal supersymmetric standard model (MSSM) contains a very large number of unknown mass parameters. Given some theoretical assumptions, and the available accelerator constraints on some of these masses, cosmology becomes a useful tool in further constraining this parameter space when the lightest supersymmetric particle (LSP) is stable [1]. As the LSP is a potential dark matter candidate, regions in parameter space can be excluded when the relic density of the LSP is excessive. Though there are large portions of the total parameter space in which the LSP in fact comprises much of the dark matter necessary to obtain closure density (or the flat rotation curves of spiral galaxies) [2, 3], direct or indirect detection of the LSP is most favorable when it is relatively light and is a mix of higgsinos and gauginos. Light particles are preferable since a fixed mass density provides a larger flux of LSPs at low masses, while lower mass sfermions (to keep the cosmological density down) make for larger elastic scattering cross-sections (see. eg. [4]).

There is however a curious but surprisingly natural possibility that all of the sfermions are in fact very heavy except for the lighter of the two stop quarks [5]. In this case, as has been recently pointed out [6], the LSP could very well be a light photino (with $m_{\tilde{\gamma}} \sim 20$ GeV) with an acceptable cosmological abundance. In what follows below, we determine the extent to which the photino must be degenerate with the light stop and examine the implications for the detection of such a light LSP. We find that despite its low mass, and the low mass of the stop, this LSP is remarkably invisible to direct detection and even to most methods of indirect detection. Even the potentially most favorable possibility, which in this case seems to be the cold annihilation to photon pairs in the galactic halo gives a signal well below background. We will also discuss the implications of this scenario on the total width of the top quark relative to the favored decay channel $t \to bW$. Light stops have also been considered recently in other contexts as well [7].

To easily see why one might expect a light stop, one needs only to examine the general form of the sfermion mass matrix [5]

$$\left(\tilde{f}_L^* \quad \tilde{f}_R^* \right) \begin{pmatrix} M_L^2 & m^2 \\ m^2 & M_R^2 \end{pmatrix} \begin{pmatrix} \tilde{f}_L \\ \tilde{f}_R \end{pmatrix} , \tag{1}$$

where $m^2 = m_f(A_f + \mu \cot\beta)$ for weak isospin $+1/2$ fermions and $m_f(A_f + \mu \tan\beta)$ for weak isospin $-1/2$. A_f is the soft supersymmetry breaking trilinear mass term, and μ is the Higgs mixing parameter; see [8] for an overview of couplings in the MSSM. This mass matrix is

easily diagonalized by writing the diagonal sfermion eigenstates as

$$\begin{aligned} \tilde{f}_1 &= \tilde{f}_L \cos\theta_f + \tilde{f}_R \sin\theta_f \ , \\ \tilde{f}_2 &= -\tilde{f}_L \sin\theta_f + \tilde{f}_R \cos\theta_f \ . \end{aligned} \tag{2}$$

With these conventions we have the mass eigenvalues

$$m_{1,2}^2 = \frac{1}{2}\left[(M_R^2 + M_L^2) \mp \sqrt{(M_R^2 - M_L^2)^2 + 4m^4}\right] \ . \tag{3}$$

Note that in the special case $M_L = M_R = M$, we have $\theta_f = \mathrm{sign}[-m^2](\pi/4)$ and $m_{1,2}^2 = M^2 \mp |m^2|$.

Now let us suppose that all of the supersymmetry mass scales (M_R, M_L, and A_f) are large ($\gtrsim$ 300 GeV) except for the gaugino masses which remain relatively light. For the SU(2) gaugino mass $M_2 \lesssim 50$ GeV and $\mu \lesssim 300$ GeV, the LSP is predominantly a photino (see eg. [9]). Note however that for large $\tan\beta(\gtrsim 3)$ or for $\mu > 0$ (unless $\tan\beta \lesssim 1.2$) much of this parameter space is excluded due to the experimental absence of light charginos. Thus for large A_f, it is quite plausible that m^2 is comparable to M_R and M_L for tops since $m_t \approx 174$ GeV [10]. As we will see, this region of parameter space is allowed, and it provides us with a relatively light cold dark matter candidate (a photino) which surprisingly would be very difficult to detect.

The light scalar top $\tilde{t}_1$ pair coupling to the Z^o is proportional to $\frac{1}{2}\cos^2\theta_t - \frac{2}{3}\sin^2\theta_W$ and vanishes for $\theta_t \simeq 0.98$ rad. For this value of θ_t, direct searches for $\tilde{t}_1$ at e^+e^- machines rely on photon mediated production. A published report [11] by the VENUS Collaboration at TRISTAN ($\sqrt{s} = 58$ GeV) excludes $\tilde{t}_1$ in the mass range between 7.6 and 28.0 GeV, for nearly all values of $m_{\tilde{\gamma}}$, except those very close to the kinematical limit for the dominant decay $\tilde{t}_1 \to c\tilde{\gamma}$ (in fact, the bound does not depend on the precise identification of the lightest neutralino). There are no published reports of similar searches at LEP, although preliminary results from the DELPHI Collaboration have appeared in conference proceedings [12]. The results allow a light scalar top below $M_Z/2$ provided that $0.9 \lesssim \theta_t \lesssim 1.1$, and that $\tilde{t}_1$ and $\tilde{\gamma}$ have masses that are not too different, which is the case we will consider here.

In order to avoid an excessive cosmological mass density, a photino with mass $m_{\tilde{\gamma}} < 35$ GeV would normally require sfermion masses $m_{\tilde{f}} \lesssim 140$ GeV. For larger sfermion masses, as we are considering here, the photino annihilation cross sections will be too small to

maintain a relic density $\Omega_{\tilde{\gamma}} h^2 \lesssim 1$. The fact that we have a light stop is of no help because annihilation into a $t, \bar{t}$ pair is kinematically forbidden. However, because photinos remain in thermal equilibrium with the stops (if they are close in mass) through the process $\tilde{\gamma} + c \leftrightarrow \tilde{t}_1$ (for light stops in the mass range considered in this paper, the radiatively induced process $\tilde{t}_1 \rightarrow \tilde{\gamma} + c$ is the dominant stop decay mode [13]), the annihilation of stops can ensure a sufficiently low relic density of photinos [6]. In what follows below, we will refine this calculation and determine the allowable photino mass as a function of the light stop mass. We will also discuss the potential (or lack thereof) for detecting this light photino. Finally, we will discuss the implications of these parameters on the top quark branching ratios.

Although the photino annihilation cross section is too small to reduce the relic abundance of photinos, the decays and inverse decays of $\tilde{t}_1 \leftrightarrow c + \tilde{\gamma}$ keep photinos in equilibrium with stops whose annihilation cross section is then sufficient for reducing the number density of both [14]. From the rate $\tilde{t}_1 \rightarrow c + \tilde{\gamma}$ [13]

$$\Gamma(\tilde{t}_1 \rightarrow c + \tilde{\gamma}) = (0.3 - 3) \times 10^{-10} m_{\tilde{t}_1} \left(1 - \frac{{m_{\tilde{\gamma}}}^2}{{m_{\tilde{t}_1}}^2}\right)^2 \tag{4}$$

it is easy to see that when compared to the expansion rate of the Universe,

$$H = \frac{8\pi}{3} N^{1/2} T^2 / M_P \tag{5}$$

where N is the number of relativistic degrees of freedom at temperature T, photinos and stops will always be in equilibrium at $T \lesssim 10^5$ GeV, unless stops and photinos are *extremely* degenerate. Therefore to determine the relic density of photinos one needs only the annihilation cross section of stops into gluons [6] (the co-annihilation of photinos and stops to charm plus glue is sufficiently suppressed that it can be ignored in this context)

$$\sigma v = \frac{14\pi\alpha_s^2}{27 {m_{\tilde{t}_1}}^2} \tag{6}$$

The Boltzmann equation for this system is most simply given in terms of the sum of the number densities $n = n_{\tilde{\gamma}} + n_{\tilde{t}_1}$ [15, 14]

$$\frac{dn}{dt} = -3Hn - \langle \sigma_{\rm eff} v \rangle \left(n^2 - {n_{\rm eq}}^2\right) \tag{7}$$

where the effective cross-section is defined to be

$$\langle \sigma_{\rm eff} v \rangle = 2\sigma v r^2 \qquad r = \frac{n_{\tilde{t}_1}}{n} \simeq \frac{3}{2} \left(\frac{m_{\tilde{t}_1}}{m_{\tilde{\gamma}}}\right)^{3/2} e^{-\Delta/x} \tag{8}$$

with $\Delta = (m_{\tilde{t}_1} - m_{\tilde{\gamma}})/m_{\tilde{\gamma}}$ and $x \equiv T/m_{\tilde{\gamma}}$.

It is relatively straightforward to integrate the Boltzmann equation (7). The result is simply

$$n = \left(\frac{4\pi^3 N}{45}\right)^{1/2} \frac{T^3}{m_{\tilde{\gamma}} M_P} \left(\int_{x_f}^{\infty} \langle \sigma_{\rm eff} v \rangle dx\right)^{-1} \tag{9}$$

where x_f corresponds to the temperature at freeze-out determined by the condition

$$\frac{.0046}{\sqrt{N}} \left(\frac{m_{\tilde{t}_1} M_P}{{m_{\tilde{\gamma}}}^2}\right) x^{-1/2} e^{-(2\Delta+1)/x} \simeq 1 \tag{10}$$

Integrating the cross-section, gives

$$\int_{x_f}^{\infty} \langle \sigma_{\rm eff} v \rangle dx \simeq 0.21 \left(\frac{\Delta m_{\tilde{t}_1}}{{m_{\tilde{\gamma}}}^3}\right) \Gamma(-1, 2\Delta/x_f) \simeq \frac{0.05}{\Delta} \left(\frac{m_{\tilde{t}_1}}{{m_{\tilde{\gamma}}}^3}\right) e^{-2\Delta/x_f} x_f^2 \left(1 - \frac{x_f}{\Delta}\right) \tag{11}$$

so that

$$\Omega_{\tilde{\gamma}} h^2 = \frac{m_{\tilde{\gamma}} n}{\rho_c} = 1.1 \times 10^{-10} \left(\frac{(m_{\tilde{\gamma}}/{\rm GeV})^3}{\Delta \Gamma(-1, 2\Delta/x_f)}\right) \left(\frac{T_{\tilde{\gamma}}}{T_\gamma}\right)^3 \sqrt{N} \tag{12}$$

The result for the relic photino density is shown in figure 1 as a function of the photino mass for different values of the stop mass. (Note that for the curve labeled 20, corresponding to the choice of the stop mass, only the portion of the curve above $\sim$ 17 Gev is allowed by VENUS results [11].) These results are qualitatively similar (and close quantitatively) to those in ref. [6]. It is clear that for a wide range of photino masses between 17 and 33 GeV, the relic density is significant so long as the stop mass is close (within 3 to 7 GeV) to that of the photino mass.

Let us now turn to the potential for detecting this light photino. Detection of the LSP can be divided among direct methods (typically involving cryogenic laboratory detectors) [4, 16] and indirect methods which make use of the annihilation of LSP's trapped in the sun or the earth [17] or annihilations in the halo of the galaxy [18]-[20]. The direct detection methods as well as the indirect methods looking for the annihilation products of trapped LSP's, depend crucially on the elastic scattering cross-section of LSP's on matter. For pure (or nearly pure) photinos, as is the case we are considering here, the elastic cross sections are suppressed by ${m_{\tilde{f}}}^4$ where $m_{\tilde{f}}$ is the mass of the heavy and potentially very heavy squarks. Since the top quark content of the proton is negligible, the light stop makes virtually no contribution to the elastic scattering cross section. Typically, for elastic scattering mediated

by sfermions, the scattering cross section leads to detection rates which are in fact marginal when $m_{\tilde{f}} \simeq 100 GeV$. For the sfermion masses we consider here, these rates would be lowered by two orders of magnitude or more.

For heavier LSP's, $m_\chi > 35$ GeV, (LSP's with larger values of M_2 and similar values of μ compared to the photinos considered here), the LSP, is no longer a pure photino, but rather becomes more of a generic mix between a gaugino and higgsino [9]. In this case, if there is in addition a sufficiently light Higgs boson, the elastic scattering may be enhanced [21]. However, for $m_{\tilde{\gamma}} \lesssim 35$ GeV, we have the curious situation where we are allowed a light photino (if the stop is nearly as light) which at the same time, if the all the other sfermions are very heavy, is essentially completely transparent to matter.

An indirect method of detection of relic photinos as a dominant component of the Galactic Halo that is also of interest is to search for a narrow line in the cosmic gamma ray spectrum at $E_\gamma = m_{\tilde{\gamma}}$, produced by the annihilation reaction $\tilde{\gamma}\tilde{\gamma} \to \gamma\gamma$ in the halo [19]-[20]. Four types of Feynman diagrams contribute to this process (together with their crossed partners) as shown in Fig. 2, where the internal lines correspond to either scalars (dashed) or fermions (full). We consider the contributions of all fermions but the top quark, for which $m_f \ll m_{\tilde{f}}$, separately. For these fermions, the dominant contribution to the annihilation amplitude is found to be of order $(eQ_f)^4 m_{\tilde{\gamma}}^2 / m_{\tilde{f}}^2$: the contribution of a heavy top quark with $m_t \simeq m_{\tilde{t}_1}$ is negligible. The observability of the gamma ray line flux under these conditions has been considered in detail in [20]: as this flux scales like ${m_{\tilde{f}}}^{-4}$, the signal dwindles into insignificance if values of $m_{\tilde{f}}$ of several hundred GeV are considered.

The contribution of the top quark to this process must however be reconsidered in the light stop scenario, $m_t \gg m_{\tilde{t}_1}$. Indeed, if all other scalar superpartners are taken to be much heavier, one could hope that top intermediate states could still lead to an observable signal. This, unfortunately, is not what happens. In the light stop scenario, the leading contribution from top intermediate states will come from diagrams 2(c) and 2(d), which together form a gauge invariant subset. Their contribution is reliably calculated in the local limit $m_{\tilde{t}_1} \ll m_t$ by first deriving the effective lagrangian $\mathcal{L}_{\rm eff}$ for $\tilde{\gamma}\tilde{\gamma} \to \tilde{t}_1 \tilde{t}_1^*$ and obtaining the $\tilde{\gamma}\tilde{\gamma} \to \gamma\gamma$ matrix element

$$\mathcal{M}(\tilde{\gamma}\tilde{\gamma} \to \gamma\gamma) = \langle \gamma\gamma | \mathcal{L}_{\rm eff} | \tilde{\gamma}\tilde{\gamma} \rangle \tag{13}$$

The effective Lagrangian for $\tilde{\gamma}\tilde{\gamma} \to \tilde{t}_1\tilde{t}_1^*$ has the leading terms:

$$\mathcal{L}_{\text{eff}} = \frac{4e^2}{9m_t}\left[\frac{m_{\tilde{\gamma}}}{m_t} - \sin 2\theta_t\right]\bar{\lambda}(x)\lambda(x)\tilde{t}_1^\dagger(x)\tilde{t}_1(x) + \frac{2e^2}{9m_t^2}\cos 2\theta_t\bar{\lambda}(x)\gamma_\mu\gamma_5\lambda(x)\tilde{J}^\mu(x) \tag{14}$$

where $\lambda(x)$ is the (Majorana) photino field and $\tilde{J}^\mu(x)$ the vector current of the scalar tops,

$$\tilde{J}^\mu(x) = \ i\ \tilde{t}_1^\dagger(x)\overleftrightarrow{\partial^\mu}\tilde{t}_1(x) \tag{15}$$

Now one can quickly see what results for the $\tilde{\gamma}\tilde{\gamma} \to \gamma\gamma$ process: the matrix element is written as,

$$\langle\gamma\gamma|\mathcal{L}_{\text{eff}}|\tilde{\gamma}\tilde{\gamma}\rangle = \frac{4e^2}{9m_t}\left[\frac{m_{\tilde{\gamma}}}{m_t} - \sin 2\theta_t\right]\langle\gamma\gamma|\tilde{t}_1^\dagger\tilde{t}_1|0\rangle\langle 0|\bar{\lambda}\lambda|\tilde{\gamma}\tilde{\gamma}\rangle + \frac{2e^2}{9m_t^2}\cos 2\theta_t\langle\gamma\gamma|\tilde{J}^\mu(x)|0\rangle\langle 0|\bar{\lambda}\gamma_\mu\gamma_5\lambda|\tilde{\gamma}\tilde{\gamma}\rangle \tag{16}$$

The second term vanishes identically by C-invariance, $\langle\gamma\gamma|\tilde{J}^\mu(x)|0\rangle \equiv 0$, since the current $\tilde{J}^\mu$ is C-odd, while the vacuum and two- photon states are C-even. In the first term, the matrix element $\langle\gamma\gamma|\tilde{t}_1^\dagger\tilde{t}_1|0\rangle$ does not vanish, and is readily evaluated: one finds,

$$\langle\gamma(k_1\gamma(k_2)|\tilde{t}_1^\dagger\tilde{t}_1|0\rangle = \frac{2e^2}{3\pi^2 m_{\tilde{t}_1}^2}(\eta^{\mu\nu}k_1\cdot k_2 - k_1^\nu k_2^\mu)\epsilon_\mu^*(k_1)\epsilon_\nu^*(k_2)I(m_{\tilde{\gamma}}/m_{\tilde{t}_1}) \tag{17}$$

where

$$I(\xi) = \int_0^1 dy\int_0^1 dx\frac{x^3y(1-y)}{1-4\xi^2x^2y(1-y)} \tag{18}$$

This integral is fairly easily expressible in closed form: we shall not bother here, because now it is the second factor of the first term that suppresses this matrix element enormously. Indeed, for non-relativistic ($\beta = v/c \simeq 10^{-3}$) relic photinos in the halo, the matrix element of the scalar density is found to be suppressed,

$$\langle 0|\bar{\lambda}\lambda|\tilde{\gamma}\tilde{\gamma}\rangle \simeq \mathrm{O}(m_{\tilde{\gamma}}\beta) \tag{19}$$

As a result, the amplitude due to stop exchange behaves parametrically as $\beta m_{\tilde{\gamma}}^3/m_{\tilde{t}_1}^2 m_t$, to be compared with the light fermion contribution discussed earlier which goes like $m_{\tilde{\gamma}}^2/m_{\tilde{f}}^2$. It is clear without any further ado (numerical factors are easily taken into account) that given the range of $m_{\tilde{t}_1}$ considered here and the necessary constraint on $m_{\tilde{\gamma}}/m_{\tilde{t}_1}$, the top contribution is always substantially less than that of the light fermions, even for the largest conceivable values of $m_{\tilde{f}}$, in which case the line signal is all but unobservable anyway, as

mentioned before. Given that the leading contribution from diagrams 2(c) and 2(d) with top intermediate states is p-wave suppressed (Eqs. 16,19), we need to discuss the contributions of diagrams 2(a) and 2(b) as well. Here, the dominant contribution to the amplitude comes from loop momenta of order m_t in the box graphs, so one cannot take the local limit before doing the loop integrals. Furthermore, in the limit $m^2_{\tilde{t}_1} \to 0$ (i.e. $m^2_{\tilde{t}_1} \ll m^2_t$), the diagrams 2(a) and 2(b) are convergent in the infrared, so unlike the case of diagrams 2(c) and 2(d), there will be no inverse powers of $m^2_{\tilde{t}_1}$. Now the leading behavior can be obtained by the informed use of dimensional analysis. Under the most favorable circumstances, the parametric dependence of the matrix element from these diagrams would be $m^3_{\tilde{\gamma}}/m^3_t$: whether or not there is also a p-wave suppression factor can only be decided by explicit calculation. If $m_{\tilde{f}} \simeq 2m_t$, say, this could be of the same order of magnitude as the $m^2_{\tilde{\gamma}}/m^2_{\tilde{f}}$ of a given light fermion: while this would be larger than the contribution from the other pair of diagrams, the gamma ray line flux would still be too low to be observable.

Thus, we come to the conclusion that in the light stop scenario, with multi-hundred GeV scalar superpartners, the $\tilde{\gamma}\tilde{\gamma} \to \gamma\gamma$ process in the Halo could not lead to an observable cosmic gamma ray line signature.

The light stop scenario has many implications in the MSSM, for example as regards the ρ parameter, $K^o - \bar{K}^o$ mixing, the rare decay $b \to s\gamma$, proton decay through dimension-5 operators in MSGUT, and generally the light sparticle spectrum. A review of these issues in [6] concludes that this scenario is perfectly allowed by known phenomenology. Here, in view of the recent announcement by the CDF collaboration [10] of evidence for the top quark we wish to consider only the implications of the light stop scenario on top decay branching ratios. In addition to the canonical decay channel

$$\Gamma(t \to bW) = \frac{G_F m_t^3}{8\sqrt{2}\pi}\left(1 - \frac{M_W^2}{m_t^2}\right)^2\left(1 + \frac{2M_W^2}{m_t^2}\right) = 1.53\mathrm{GeV} \tag{20}$$

there is the possibility for decays into photinos and stops with a rate given by [5]

$$\Gamma(t \to \tilde{t}_1\tilde{\gamma}) = \frac{\alpha}{9m_t^3}(m_t^2 + m^2_{\tilde{\gamma}} - m^2_{\tilde{t}_1} - 2m_t m_{\tilde{\gamma}}\sin 2\theta_t)\lambda^{1/2}(m_t^2, m^2_{\tilde{\gamma}}, m^2_{\tilde{t}_1})\theta(m_t - m_{\tilde{\gamma}} - m_{\tilde{t}_1}) \tag{21}$$

For stops with $m_{\tilde{t}_1} \lesssim 32 GeV$, (assuming the GUT relation between the photino and gluino masses) there is also the possibility that tops decay to stops and gluinos with a rate given by eq. (21) with the substitution $4\alpha/9 \to 4\alpha_s/3$ and $m_{\tilde{\gamma}} \to m_{\tilde{g}}$. (We implicitly assume that

the decay $t \to bH^+$, a possibility in the MSSM, is absent due to kinematics, i.e. we assume that the H^+ is heavy.) The branching ratio for the $t \to bW$ decay is shown in figure 3 as the area between the two curves. In computing these curves, we have assumed the GUT relation between the gaugino masses so that $m_{\tilde{g}} = 5.76 m_{\tilde{\gamma}}$ and we have taken the value of θ_t to be that for which the coupling of the light stop to the Z^o vanishes. The upper curve corresponds to the branching ratio when $(m_{\tilde{t}_1} - m_{\tilde{\gamma}}) = 2 - 3$ GeV (the minimum allowed from figure 1 for this range of stop masses) and the lower curve corresponds to the branching ratio when $(m_{\tilde{t}_1} - m_{\tilde{\gamma}}) = 3 - 7$ GeV (the maximum allowed from figure 1 for this range of stop masses). As one can see, for $m_{\tilde{t}_1} \gtrsim 30$ GeV, the only additional channel is the decay to stops and photinos which makes a 5 % contribution to the total width. At lower values of $m_{\tilde{t}_1}$ the contribution from $t \to \tilde{t}_1 \tilde{g}$ becomes significant. The reduced $t \to bW$ branching ratio would then worsen the discrepancy [10] between the CDF top quark signal and the expected QCD rate for $p\bar{p} \to t\bar{t}X$ production. Thus, it would appear that such considerations would disfavor the low range of mass values $m_{\tilde{t}_1} \lesssim 30$ GeV. However, it should be borne in mind that gluinos in the corresponding mass range, up to ~ 145 GeV will be pair produced in $p\bar{p}$ collisions, with a sizable branching ratio for the decay $\tilde{g} \to \tilde{t}_1 \bar{t}(\text{virtual}) \to (c\tilde{\gamma})(\bar{b}W)$, providing a new source of events with b jets and W's which should be taken into account in data analysis.

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Figure Captions

Figure 1: The relic density of photinos as a function of the photino mass. The different curves are labeled by the stop quark mass in GeV.

Figure 2: Typical box diagrams contributing to the transition amplitude for the process $\tilde{\gamma}\tilde{\gamma} \to \gamma\gamma$. The corresponding diagrams with photon lines interchanged are not shown.

Figure 3: The branching ratio for $t \to bW$ as a function of the stop mass. The two curves encompass the allowed range of photino masses which give a cosmologically allowed and interesting relic density from figure 1.